

Radiative decays $J/\Psi \rightarrow \eta^{(\prime)}\gamma$ in perturbative QCD

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Abstract

With the recent investigations of $g^*g^* - \eta^{(\prime)}$ transition form factor and $\eta - \eta'$ mixing scheme, we present an updated study of the radiative decays $J/\Psi \rightarrow \eta^{(\prime)}\gamma$ in perturbative QCD. The decays are taken as a test ground for the $g^*g^* - \eta^{(\prime)}$ transition form factors and the $\eta - \eta'$ mixing scheme. The form factors are found to be working for glunic $\eta^{(\prime)}$ productions and the mixing angle is constrained to be $\phi = 35.1^\circ \pm 0.8^\circ$.

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As it is well known that the heavy quarkonium decays to light mesons have played very important role in testing and understanding QCD from the very beginning. The decays $J/\Psi \rightarrow \eta^{(\prime)}\gamma$ are of great interests since they are closely related to the issues of $\eta' - \eta$ mixing and $g^*g^* - \eta^{(\prime)}$ transition form factors, which are very important ingredients for understanding many interesting hadronic phenomena of η and η' productions. For example, it would be very useful for explaining the large branching ratio of strong penguin dominated decay $B \rightarrow K\eta'$ [1, 2, 3].

In the literature, studies of the decays $J/\Psi \rightarrow \eta^{(\prime)}\gamma$ are different each other from the treatments of formatting gluons to $\eta^{(\prime)}$, namely, direct nonperturbative $gg - \eta^{(\prime)}$ coupling through strong anomaly[4] or two off-shell gluons coupled to $\eta^{(\prime)}$ through quark loop[5]. In this letter, we will take the second approach which had been pioneered systematically within perturbative QCD by Körner, Kühn, Krammer and Schneider(KKKS) [5] years ago. In Ref.[5], the non-relativistic quark model and the weak-binding approximation were used for both heavy and light mesons, and systematic helicity projectors were constructed to reduce loop integrations. In this work, we follow their approach. However, two improvements are included:

- $g^*g^* - \eta^{(\prime)}$ couplings are improved to be relativistic transition form factors as advocated in Ref.[6, 7, 8] in stead of non-relativistic modelling.
- The $\eta' - \eta$ mixing scheme is also updated to the Feldmann-Kroll-Stech(FKS) mixing scheme[9].

In perturbative QCD approach, the decays are depicted by the Feynman diagrams in Fig.1. To calculate the amplitudes for the decays, we need to know how to deal with the dynamics of bound states. Generally, factorization are employed. Soft nonperturbative QCD bound state dynamics are factorized to the decay constants and the wave functions of J/Ψ and $\eta^{(\prime)}$ which will convolute with the hard kernel induced the decay. We shall use the non-relativistic approximation for the heavy J/Ψ , but not for the light mesons η' and η . Although a rigorous theory from first principles for the light bound-states are still missing, some effective approaches are in progress. In recent years, it has been realized that a proper treatment of the $\eta - \eta'$ system requires a sharp distinction between the mixing states and the mixing properties of the decay constants[9]. Taking strange-nonstrange flavor basis for the $\eta - \eta'$ system and the mixing of the decay constants following the same pattern of the state mixing, FKS have found a dramatic

simplification. They also have tested their mixing scheme against experiment and determined corrections to the first order values of the basic parameters from phenomenology.

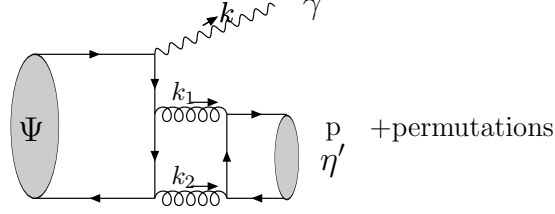


Figure 1: Lowest order QCD diagrams for $J/\Psi \rightarrow \eta^{(l)}\gamma$ decays.

In FKS mixing scheme the parton Fock state decomposition can be expressed as

$$\begin{aligned} |\eta\rangle &= \cos\phi |\eta_q\rangle - \sin\phi |\eta_s\rangle, \\ |\eta'\rangle &= \sin\phi |\eta_q\rangle + \cos\phi |\eta_s\rangle \end{aligned} \quad (1)$$

where ϕ is the mixing angle, $|\eta_q\rangle \sim f_q \Phi(x, \mu) |u\bar{u} + d\bar{d}\rangle / \sqrt{2}$ and $|\eta_s\rangle \sim f_s \Phi(x, \mu) |s\bar{s}\rangle$. The decay constants f_q , f_s and the mixing angle ϕ have been constrained from the available experimental data, $f_q = (1.07 \pm 0.02)f_\pi$, $f_s = (1.34 \pm 0.06)f_\pi$, $\phi = 39.3^\circ \pm 1.0^\circ$ [9].

Already in Ref.[10], Baier and Grozin have derived evolution equations for the distribution functions $\Phi(x, \mu)$ to the first order of α_s , which eigenfunctions are found to be

$$\Phi(x, \mu) = 6x(1-x) \left(1 + \sum_{n=2,4,\dots} B_n(\mu) C^{\frac{3}{2}}(2x-1) \right). \quad (2)$$

In the limit $\mu \rightarrow \infty$, the coefficients B_n evolve to zero and $\Phi(x, \mu)$ turns out to be $\phi_{AS} = 6x(1-x)$. When evolution equations run down to low energy scale, its quark contents mixed with glunic states. However, the gluon content enters the $\eta^{(l)}$ wave function from next-to-leading order. This observation encourages the calculations of the $g^*g^* - \eta^{(l)}$ transition form factors similarly to the well known $\gamma^* - \pi$ transition form factor at leading order, which read [6, 7, 8]

$$\begin{aligned} \mathcal{M}_{\mu\nu} &= \langle g_a^* g_b^* | \eta^{(l)} \rangle = -4\pi\alpha_s \delta_{ab} i\epsilon_{\mu\nu\alpha\beta} Q_1^\alpha Q_2^\beta F_{g^*g^*-\eta^{(l)}}(Q_1^2, Q_2^2), \\ F_{g^*g^*-\eta^{(l)}}(Q_1^2, Q_2^2) &= \frac{1}{2N_c} f_{\eta^{(l)}} \int_0^1 dx \frac{\phi_{\eta^{(l)}}(x, \mu)}{\bar{x}Q_1^2 + xQ_2^2 - x\bar{x}m_{\eta^{(l)}}^2 + i\epsilon} + (x \rightarrow \bar{x}), \end{aligned} \quad (3)$$

where $\bar{x} = 1-x$, $f_{\eta'} = \sqrt{2}f_q \sin\phi + f_s \cos\phi$ and $f_\eta = \sqrt{2}f_q \cos\phi - f_s \sin\phi$. To the accuracy of this paper, $\phi_{\eta^{(l)}}(x, \mu)$ is taken to be the leading twist distribution functions(DAs) $\phi_{\eta^{(l)}}^{AS}(x) = 6x(1-x)$.

Using $g^*g^* - \eta^{(\prime)}$ in Eq.3 and following the procedure developed in Ref.[5, 11], it is straightforward to evaluate the amplitudes for the decays as depicted by Feynman diagrams in Fig.1. We get

$$\Gamma(V \rightarrow \eta'\gamma) = \frac{1}{6} \left(\frac{2}{3}\right)^2 e_Q^2 \alpha_s^4(M_V) \alpha_e \frac{f_V^2 f_{\eta'}^2}{M_V^3} (1 - z^2) |H(z)|^2, \quad (4)$$

where $z = m_{\eta'}/M_V$, e_Q is the heavy quark electric charge and $\frac{2}{3}$ is the color factor. The dimensionless scalar function $H(z)$ containing loop integrals is given by

$$H(z) = \frac{M_V^2}{2p \cdot k} \frac{1}{16} \frac{1}{i\pi^2} \int_0^1 du \phi_{\eta'}^{AS}(u) \int d^4q \frac{k_1 \cdot k_2 (p \cdot k q^2 - q \cdot k q \cdot p)}{D_1 D_2 k_1^2 k_2^2 (\bar{u} k_1^2 + u k_2^2 - u \bar{u} m_{\eta'}^2)}, \quad (5)$$

where $D_1 = -k_1 \cdot (k + k_2)$, $D_2 = -k_2 \cdot (k + k_1)$, $q = k_1 - k_2$, and $p = k_1 + k_2$.

Obviously in Eq.5, the $k_1 \cdot k_2$ numerator would cancel the η' form factor if it is taken be $\sim 1/k_1 \cdot k_2$, and the hard scattering kernel would not convolute with the distribution functions of η' .

With the help of the algebraic identities

$$q^2 = \frac{2}{M_V^2 + m_{\eta'}^2} [m_{\eta'}^2 (D_1 + D_2) + M_V^2 (k_1^2 + k_2^2)], \quad (6)$$

$$q \cdot p = k_1^2 + k_2^2, \quad k_1 \cdot k_2 = \frac{1}{2} (p^2 - k_1^2 - k_2^2) = -\frac{1}{2} (p \cdot k + D_1 + D_2), \quad (7)$$

the integrand in $H(z)$ can be decomposed into a sum of four, three and two-points functions which is presented in appendix A. In the calculation of the loop integrals, we have used dimensional regularization scheme and the methods developed in Ref.[12]

For numerical results for the decays, we use $\Gamma_{tot.}(J/\Psi) = (87 \pm 5) \text{Kev}$ [13], $f_{J/\Psi} = 400 \text{Mev}$ and $\alpha_s(M_{J/\Psi}) = 0.2557$ [14]. We get

$$\mathcal{B}^{th}(J/\Psi \rightarrow \eta'\gamma) = 3.9 \times 10^{-3}, \quad \left(\mathcal{B}^{exp}(J/\Psi \rightarrow \eta'\gamma) = (4.3 \pm 0.3) \times 10^{-3}, PDG[13] \right) \quad (8)$$

$$\mathcal{B}^{th}(J/\Psi \rightarrow \eta\gamma) = 3.5 \times 10^{-4}, \quad \left(\mathcal{B}^{exp}(J/\Psi \rightarrow \eta\gamma) = (8.6 \pm 0.8) \times 10^{-4}, PDG[13] \right). \quad (9)$$

While $\mathcal{B}^{th}(J/\Psi \rightarrow \eta'\gamma)$ agrees with experiment, $\mathcal{B}^{th}(J/\Psi \rightarrow \eta\gamma)$ turns out to be too small. From the mixing scheme, it is easy to see that $\mathcal{B}^{th}(J/\Psi \rightarrow \eta'\gamma)$ is *insensitive* to the mixing angle ϕ when ϕ is about 35° , but $\mathcal{B}^{th}(J/\Psi \rightarrow \eta\gamma)$ is very *sensitive* to ϕ . Take $\phi = 35.3^\circ$ fitted from $\eta' \rightarrow \rho\gamma$ and $\rho \rightarrow \eta\gamma$ [9], we find

$$\mathcal{B}^{th}(J/\Psi \rightarrow \eta'\gamma) = 3.75 \times 10^{-3}, \quad \mathcal{B}^{th}(J/\Psi \rightarrow \eta\gamma) = 7.3 \times 10^{-4}, \quad (10)$$

which agree with experimental results quite well. However, if we take $\alpha_s(\mu) = \alpha_s m_c$, the results turn out to overshoot their experimental data.

The most theoretical uncertainty may arise from the energy scale choice in $\alpha_s(\mu)$. Because our calculation is performed at the lowest order in QCD and there is no UV divergence in the loop diagram induced the decay, we don't have strong argument to choose a scale, as in usual case, to minimize the higher order corrections by setting logarithm to zero. Naively, the scale could be chosen from m_c to $m_{J/\Psi}$. To reduce the scale dependence, we relate $\mathcal{B}(J/\Psi \rightarrow \eta^{(\prime)}\gamma)$ to $\mathcal{B}(J/\Psi \rightarrow ggg)$

$$\mathcal{B}(J/\Psi \rightarrow \eta^{(\prime)}\gamma) = \frac{\Gamma(J/\Psi \rightarrow \eta^{(\prime)}\gamma)}{\Gamma(J/\Psi \rightarrow ggg)} \mathcal{B}(J/\Psi \rightarrow ggg). \quad (11)$$

With the help of the known results[15]

$$\frac{\Gamma(V \rightarrow ggg)}{\Gamma(V \rightarrow \mu^+\mu^-)} = \frac{10(\pi^2 - 9)}{81\pi e_Q^2} \frac{\alpha_s^3(M)}{\alpha_e^2} \left\{ 1 + \frac{\alpha_s(M)}{\pi} \left[-19.4 + \frac{3}{2}\beta_0 \left(1.16 + \ln \left(\frac{2M}{M_V} \right) \right) \right] \right\}, \quad (12)$$

we can get

$$\mathcal{B}(J/\Psi \rightarrow \eta^{(\prime)}\gamma) = \frac{9}{20(\pi^2 - 9)} \frac{e_Q^2}{M_V^2} \frac{\alpha_s(M) \alpha_e f_{\eta^{(\prime)}}^2 (1 - z^2) |H(z)|^2}{1 + \frac{\alpha_s(M)}{\pi} \left[-19.4 + \frac{3}{2}\beta_0 \left(1.16 + \ln \left(\frac{2M}{M_V} \right) \right) \right]} \mathcal{B}(J/\Psi \rightarrow ggg). \quad (13)$$

We will use the following relation and experimental data[13] for our numerical results

$$\begin{aligned} \mathcal{B}(J/\Psi \rightarrow ggg) &= \mathcal{B}(J/\Psi \rightarrow \text{hadrons}) - \mathcal{B}(J/\Psi \rightarrow \text{virtual}\gamma \rightarrow \text{hadrons}) \\ &= (0.877 \pm 0.005) - (0.17 \pm 0.02) \\ &= 0.707 \pm 0.025. \end{aligned} \quad (14)$$

Taking $\phi = 35.3^\circ$ and $\alpha_s(M) = \alpha_s(m_c)$, we obtain

$$\mathcal{B}^{th}(J/\Psi \rightarrow \eta'\gamma) = 4.17 \times 10^{-3}, \quad \mathcal{B}^{th}(J/\Psi \rightarrow \eta\gamma) = 8.16 \times 10^{-4}, \quad (15)$$

which agree with the experiment data.

In Fig.2, we display the ratio $\mathcal{R}_{J/\Psi} = \mathcal{B}(J/\Psi \rightarrow \eta'\gamma)/\mathcal{B}(J/\Psi \rightarrow \eta\gamma)$ as a function of ϕ , in which we expect that the relativistic and the higher order QCD corrections may be concealed to large extent, so that the ratio could be predicted much more reliable than the two decay rates respectively. Comparing our results with the experimental measurement $R_{J/\Psi} = 5.0 \pm 0.6$ [13]

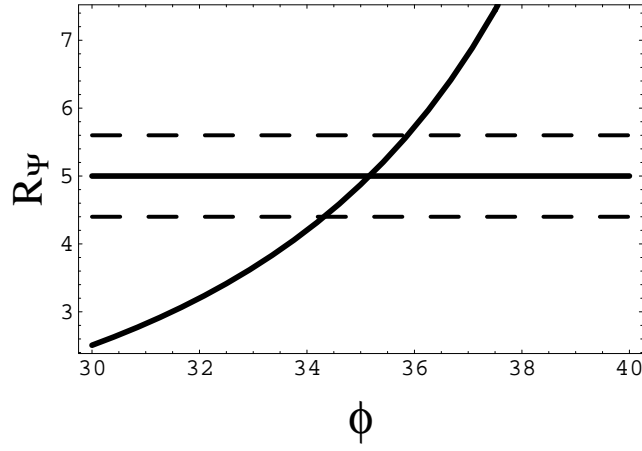


Figure 2: The ratio $\mathcal{R}_{J/\Psi}$ is shown by a solid curve as a function of ϕ (in degree). The experimental data are shown by horizontal lines. The thicker solid horizontal line is its center value, thin horizontal dash lines are its error bars.

as displayed by horizontal lines in Fig.2, we find $\phi = 35.1^\circ \pm 0.8^\circ$ which is different from the value $\phi = 39.0^\circ \pm 1.6^\circ$ by 2σ also determined from $J/\Psi \rightarrow \eta(\eta')\gamma$ [9] by using QCD anomaly dominance mechanism formula. To make clear the origin of the discrepancy between the two different determination of mixing angle ϕ , we recapitulate the key formula from the well known work of Novikov *et al.* [4]

$$R_{J/\Psi} = \left| \frac{\langle 0 | G\tilde{G} | \eta' \rangle}{\langle 0 | G\tilde{G} | \eta \rangle} \right|^2 \left(\frac{p_{\eta'}}{p_\eta} \right)^3. \quad (16)$$

This formula is frequently employed to determine $\eta - \eta'$ mixing angles in the literature. Technologically, the strong anomaly dominance is equivalent to the dominance of the ground state and the neglect of continuum contribution to dispersion relations as shown in Ref.[4, 16]. So far, considering the experimental and the theoretical uncertainties, the difference between the predictions for $\mathcal{R}_{J/\Psi}$ by the two mechanism is still marginal. Although we have improved $g^*g^* - \eta^{(\prime)}$ couplings in Ref.[5] from non-relativistic to relativistic, there is still large room for theoretical improvements which is very worth for further studying. We also note the CLEO/CESR-c project is going, where about one billion Ψ events would be produced. The refined measurements of these decays to be performed at CLEO-c will deepen our understanding of the two $\eta^{(\prime)}$ production mechanisms.

In this letter, we have studied the radiative decays $J/\Psi \rightarrow \eta'(\eta)\gamma$ in perturbative QCD. The relativistic $g^*g^* - \eta^{(\prime)}$ transition form factors have been tested to be working for η' production.

The mixing angle in FKS scheme is constrained to be $\phi = 35.1^\circ \pm 0.8^\circ$. This study encourages further applications of the form factor for $\eta^{(\prime)}$ production in hard processes. It is also very helpful for understanding the abnormal large η' yields in B meson decays, which have caught many theoretical attentions[17, 18] recently.

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Appendix A

In the evolution of the amplitudes for $J/\Psi \rightarrow \gamma\eta^{(\prime)}$, we encounter the loop integral in Eq.5 which can be expanded in terms four, three and two points functions

$$H(z) = \frac{1}{16} \frac{1}{1-z^2} \int du \phi_{AS}(u) \left[\frac{1-z^2}{2(1+z^2)} 4^4 \left(m^4 D_0^a(u, z) - \frac{1}{2} (1-z^2) m_V^4 D_0^b(u, z) \right) \right. \\ \left. - \frac{1}{2} (1-uz^2) 4^3 C_0^b(u, z) - \frac{1}{2} (1-2z^2+uz^2) 4^3 C_0^a(u, z) \right. \\ \left. + \frac{1}{2u} 4^2 (B_0^a(u, z) - B_0^b(u, z) - B_0^c(u, z) + B_0^d(u, z)) \right], \quad (17)$$

with the following functions

$$D_0^a(u, z) = \frac{1}{8m_V^4(1-u)uz^2(1-z^2)} \left[Sp \left(1 - \frac{1-\bar{u}z^2}{\bar{u}(1-z^2)} \right) + 2\pi i \ln \left(1 - \frac{1-\bar{u}z^2}{\bar{u}(1-z^2)} - i\epsilon \right) \right. \\ + Sp \left(1 - \frac{1-\bar{u}z^2}{1-(1-2u)z^2} \right) + 2\pi i \ln \left(1 - \frac{1-\bar{u}z^2}{1-(1-2u)z^2} - i\epsilon \right) - Sp \left(1 - \frac{\bar{u}-(1-2u)z^2}{\bar{u}(1-z^2)} \right) \\ \left. + \left(2\pi i + \ln \left(\frac{\bar{u}-(1-2u)z^2}{1-\bar{u}z^2} \right) \right) \left(\pi i + \ln \left(\frac{(1-z^2)(\bar{u}-(1-3u+2u^2)z^2)}{uz^2} \right) \right) \right], \quad (18)$$

$$D_0^b(u, z) = \frac{1}{4m_V^4(1-(1-2u)z^2)u(1-z^2)} \left[2Sp \left(-\frac{1-z^2}{uz^2} \right) - 2Sp \left(-\frac{(1-2u)(1-z^2)}{u} \right) \right. \\ + Sp \left(-\frac{1-z^2}{z^2(\bar{u}-(1-3u+2u^2)z^2)} \right) - Sp \left(-\frac{(1-2u)^2z^2(1-z^2)}{\bar{u}-(1-3u+2u^2)z^2} \right) \\ \left. + \ln \left(\frac{1-\bar{u}z^2}{z^2(\bar{u}-(1-2u)z^2)} \right) \ln \left(\frac{(1-z^2)(\bar{u}-(1-3u+2u^2)z^2)}{uz^2} + i\pi \right) \right], \quad (19)$$

$$C_0^b(u, z) = -\frac{1}{m_V^2} \int_0^1 dy \frac{1}{4u^2z^2 - y^2(1 - 2z^2) - 2uy(1 - 3z^2) - i\epsilon} \ln \left(\frac{y(2(1 - 2z^2) - y(1 - 2z^2) - 2u(1 - 3z^2)) - i\epsilon}{2(y(1 - z^2) - 2u^2z^2) - i\epsilon} \right), \quad (20)$$

$$C_0^a(u, z) = -\frac{1}{m_V^2} \int_0^1 dy \frac{1}{y^2 + 2y\bar{u}(1 - z^2) + 4\bar{u}z^2 - i\epsilon} \ln \left(\frac{-y(y + 2\bar{u}(1 - 2z^2) + i\epsilon)}{4\bar{u}^2z^2 + i\epsilon} \right), \quad (21)$$

$$B_0^a(u, z) = \frac{2}{\epsilon} - \gamma_E + \ln 4\pi + \ln \mu^2 + 2 - \ln m_V^2 - \left(1 - \frac{1}{(1 - 2u)(1 - 2uz^2) + i\epsilon} \right) \ln \left(1 - (1 - 2u)(1 - 2uz^2) + i\epsilon \right), \quad (22)$$

$$B_0^b(u, z) = \frac{2}{\epsilon} - \gamma_E + \ln 4\pi + \ln \mu^2 + 2 - \ln m_V^2, \quad (23)$$

$$B_0^c(u, z) = \frac{2}{\epsilon} - \gamma_E + \ln 4\pi + \ln \mu^2 + 2 - \ln m_V^2 - \left(1 - \frac{1}{(1 - 2u)(2\bar{u}z^2 - 1) + i\epsilon} \right) \ln \left[1 - (1 - 2u)(2\bar{u}z^2 - 1) + i\epsilon \right], \quad (24)$$

$$B_0^d(u, z) = \frac{2}{\epsilon} - \gamma_E + \ln 4\pi + \ln \mu^2 + 2 - \ln m_V^2 - \frac{2(1 - z^2)}{1 - 2z^2} \ln[2(1 - z^2)] + \ln(1 - 2z^2) \quad (25)$$

Where $Sp(x) = Li_2(x)$ is the Spence function and $\bar{u} = 1 - u$.

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